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TeV scale leptogenesis with heavy neutrinos

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Following a baryogenesis scenario proposed by Lazarides, Panagiotakopoulos, and Shafi, we show how the observed baryon asymmetry can be explained via resonant leptogenesis in a class of supersymmetric models with an intermediate mass scale $M_I \lesssim 10^9$ GeV. It involves the out of equilibrium decay of heavy ($\lesssim M_I$) right handed neutrinos at a temperature close to the TeV supersymmetry breaking scale. Such models can also resolve the MSSM μ problem.

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A large class of supersymmetric models possess D and F-flat directions which can have important cosmological consequences. A particularly interesting set belongs to extensions of the minimal supersymmetric standard model (MSSM) and contains one or more intermediate to superheavy scales that arise from an interplay of a TeV scale from supersymmetry breaking and higher order (nonrenormalizable) terms suppressed by some cutoff scale M_*. Such models possess the following features that were discussed quite some time ago [1–7], especially during the era of superstring inspired models:

(1) In the context of the early Universe the associated phase transition takes place at a temperature close to TeV, the supersymmetry breaking scale, even though the gauge symmetry breaking scale is of intermediate size or higher [1–3].

(2) The universe experiences a modest amount ($\sim$10 or so e-foldings) of inflation before the phase transition takes place [2–4,7]. This is now usually referred to as thermal inflation [8].

(3) An appreciable amount of entropy generation occurs at the end of inflation, and this could be exploited to dilute potentially troublesome relics such as superheavy magnetic monopoles [5].

(4) The flip side of point (3) is that either a preexisting baryon (or lepton) asymmetry should be sufficiently large to overcome the entropy onslaught, or a mechanism is in place to produce the asymmetry once the phase transition is completed. The latter case requires a final temperature of the radiation dominated Universe (T_f) in excess of an MeV (or so) to preserve hot big bang nucleosynthesis, and this sets an upper bound on the intermediate scale M_I of around $10^{15} - 10^{16}$ GeV [3].

In Refs. [2,3] a new mechanism for generating the baryon asymmetry was proposed, relying on the out-of-equilibrium decay of heavy (intermediate scale) particles at a temperature close to the TeV scale. The novel feature here is that the decaying particles acquire mass through their coupling to the scalar field that is undergoing the phase transition. Thus, the phase transition and generation of baryon asymmetry occur more or less simultaneously. As noted in Ref. [2] the gravitino problem is neatly avoided in these models.

The main purpose of this paper is to show how the scenario of Refs. [2,3] can be adapted to generate an initial lepton asymmetry, part of which is subsequently transformed to the observed baryon asymmetry [9] through electroweak sphaleron mediated transitions [10]. We invoke resonant leptogenesis [11,12] to generate the required large initial asymmetry, before its dilution from entropy production.

The scenario we have in mind naturally arises in models based on subgroups of supersymmetric $SO(10)$ such as $H_1 = SU(2)_L \times U(1)_R \times U(1)_{B-L}$ or $H_2 = SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [13]. The Higgs field ϕ, whose vacuum expectation value $\langle \phi \rangle \equiv M_I$ breaks $H_{1,2}$ to $SU(2)_L \times U(1)_Y$, should also provide masses comparable to M_I to the right-handed neutrinos. [For H_2, if ϕ belongs to the representation $(1,2)_1$, where the subscript labels the $B-L$ charge, then dimension five operators will generate masses for the right-handed neutrinos that are suppressed by the cutoff scale. However, if ϕ belongs to the representation $(1,3)_2$ of H_2, the right-handed neutrinos can acquire masses comparable to M_I. We will assume the latter case.] The renormalizable part of the superpotential contains, among others, the following terms:

$$W_R \supset f_{ij} \phi N_i N_j + h_{i\alpha} N_i L_\alpha H_u , \tag{1}$$

where N_i denote the three right-handed neutrino superfields, L_α denote the three lepton superfields, H_u is the MSSM doublet vacuum expectation value (VEV) that contributes to the neutrino Dirac mass and, unless otherwise stated, the dimensionless coefficients f_{ij} are of order unity. The Yukawa couplings $h_{i\alpha}$ should be suitably chosen to reproduce the neutrino oscillation parameters.

In order to generate an intermediate scale VEV for ϕ, the superpotential should not contain terms such as $\phi\overline{\phi}$ (the conjugate superfield $\overline{\phi}$ is present to ensure that supersymmetry is not broken at the intermediate scale). Furthermore, quartic terms consisting of the scalar component of ϕ and with coefficients of order unity must be absent from the potential. This is ensured by the gauge symmetry which forbids a cubic

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term for ϕ in the superpotential. The two most relevant dimension four (nonrenormalizable) terms in the superpotential are

$$W_{\rm NR} \supset \frac{\lambda}{M_*}(\phi\bar{\phi})^2 + \frac{\beta}{M_*}\phi\bar{\phi}H_u H_d, \tag{2}$$

where M_* denotes the cutoff scale and λ and β are dimensionless coefficients. The term proportional to β is needed, as we will see, to ensure that the final temperature after completion of the phase transition is of order $10^2 - 10^3$ GeV, so that the electroweak sphalerons can partially convert the lepton asymmetry to the observed baryon asymmetry.

The superpotential terms in Eqs. (1) and (2) are easily realized by supplementing the gauge symmetries $H_{1,2}$ with suitable additional symmetries. For instance, in the H_1 case, a discrete symmetry Z_2 under which only N_i, $\bar{\phi}$, H_d, and L_α change sign is adequate. In the H_2 case, we can use a Z_4 symmetry with the following transformations: $(H_u, H_d) \to i(H_u, H_d)$, $N_i \to -iN_i$, $\phi \to -\phi$, with $\bar{\phi}, L_\alpha$ left unchanged. Such discrete symmetries may lead to the production of domain walls which, in principle, can be problematic. A resolution of the domain wall problem in this class of models has been extensively discussed in Ref. [4].

We see from Eq. (2) that the combinations $H_u H_d$ and $\phi\bar{\phi}$ transform identically under any additional symmetries. Since $\phi\bar{\phi}$ is absent from the superpotential in order to generate a flat direction, we are led to conclude that the 'bare' MSSM μ term must also be absent. Thus, we have a nice mechanism for resolving the MSSM μ problem. The induced μ term $\beta M_I^2/M_*$ is of TeV scale as desired. (A resolution of the μ problem in this class of models has previously been discussed in Ref. [1], as well as in the first paper in Ref. [14].)

Following common practice, we use ϕ to also denote the scalar component of the superfield ϕ. Assuming that ϕ (sometimes referred to as a "flaton" [6]) has sufficiently strong Yukawa couplings [Eq. (1)] which can change the sign of its positive supersymmetry breaking mass squared term generated at some superheavy scale $\gg M_I$, and taking a D-flat direction where $\langle\phi\rangle = \langle\bar{\phi}\rangle^\dagger$, the zero-temperature effective potential of ϕ is [2,3]

$$V_0(\phi) = \mu_0^4 - M_s^2|\phi|^2 + \frac{8\lambda^2}{M_*^2}|\phi|^6. \tag{3}$$

Here $\mu_0^4 = (\frac{2}{3}M_s M_I)^2$ is included to ensure that at the minimum $\langle\phi\rangle = M_I = (\lambda^{-1}M_s M_*/2\sqrt{6})^{1/2}$, $V(M_I)=0$, and M_s ($\sim$ TeV) refers to the supersymmetry breaking scale.

For nonzero temperature the effective potential acquires an additional contribution, given by [15]

$$V_T(\phi) = \left(\frac{T^4}{2\pi^2}\right)\sum_i (-1)^F \int_0^\infty dx\, x^2 \times \ln\left(1-(-1)^F \exp\left\{-\left[x^2+\frac{M_i^2(\phi)}{T^2}\right]\right\}\right), \tag{4}$$

where the sum is over all helicity states, $(-1)^F$ is ± 1 for bosonic and fermionic states, respectively, and M_i is the field-dependent mass of the ith state. For $\phi \ll$ T Eq. (4) yields a temperature-dependent mass term $\sigma \mathrm{T}^2|\phi|^2$, where $\sigma \sim 0.2$ for $f_{ij} \sim 1$. Hence the potential

$$V(\phi) = \mu_0^4 + (-M_s^2 + \sigma \mathrm{T}^2)|\phi|^2 + \frac{8\lambda^2}{M_*^2}|\phi|^6 \tag{5}$$

has a minimum $V(\phi) = \mu_0^4$ at $\phi=0$ for $T > T_c = M_s/\sigma^{1/2}$. For $\phi > T$, the temperature-dependent term is exponentially suppressed and $V(\phi)$ develops another minimum at $\phi = M_I$ for $T \lesssim M_I$. $\phi = 0$ is the absolute minimum for $\mu_0 \lesssim T \lesssim M_I$ since the symmetric phase ($\phi=0$) has more massless degrees of freedom and the radiation energy density dominates over the false vacuum energy density μ_0^4. For $T \lesssim \mu_0$ the broken phase ($\phi = M_I$) becomes the absolute minimum of the potential, with $V(M_I)=0$. [The recently measured vacuum energy density of order $(10^{-3}\ \mathrm{eV})^4$ is negligible for our purposes.]

The universe remains at $\phi=0$ for $T > T_c$ and, for $M_I \sim 10^8$ GeV, experiences roughly $\ln(\mu_0/T_c) \sim 6$ e-foldings of inflation due to the false vacuum energy density μ_0^4 [2–4]. During this phase the right-handed neutrinos N_i are in thermal abundance

$$\frac{n_{N_i}}{s} \simeq \frac{n_{N_i}^{eq}}{s} = \frac{45\zeta(3)}{2\pi^4}\frac{1}{g_*}\left(\frac{3}{4}g_{N_i} + g_{\tilde{N}_i}\right) \simeq \frac{1}{300}, \tag{6}$$

where g_* counts the effectively massless degrees of freedom and $(g_{\tilde{N}_i})g_{N_i}$ counts the degrees of freedom of (s)neutrinos. When the temperature reaches T_c, the minimum (and the associated barrier) at $\phi=0$ disappears, and ϕ starts to roll down towards the minimum at $\phi = M_I$. The classical evolution of ϕ field is governed by the equation

$$\ddot{\phi} + 3H\dot{\phi} = -\frac{dV}{d\phi}. \tag{7}$$

For $T < \phi < M_I$ the temperature-dependent mass term and the term proportional to $|\phi|^6$ can be ignored. Also, with the Hubble constant $H = \mu_0^2/\sqrt{3}M_P \ll M_s$ (where $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass), Eq. (7) yields

$$\ddot{\phi} \simeq M_s^2 \phi, \tag{8}$$

so that

$$\phi(\delta t) \simeq T_c \exp(M_s \delta t). \tag{9}$$

From Eq. (9), it takes the flaton $\delta t \simeq \ln[M_I/T_c]/M_s \sim 10 M_s^{-1}$ to roll down to its minimum at M_I [3,4].

As the flaton rolls down, the right-handed neutrinos pick up a mass proportional to $\langle\phi\rangle$ and can decay out of equilibrium via the couplings $h_{i\alpha}N_i L_\alpha H_u$. The decay width is $\Gamma_{N_i} \simeq \Sigma_\alpha |h_{i\alpha}|^2 M_{N_i}/8\pi$, where M_{N_i} is the mass of the right-handed neutrinos when they decay.

We show later that the resulting lepton asymmetry can account for the present baryon asymmetry of the universe for $M_s \ll M_I \lesssim (M_s/\mathrm{TeV}) \times 10^8$ GeV. (We will assume throughout that M_I and $M_{N_i} \gg M_s$, otherwise thermal effects and direct flaton decay into neutrinos would modify our discussion.) Since $M_{N_i} \lesssim M_I$, the light neutrino masses ($m_{\nu_i} \lesssim 0.1$ eV [16]) require that the Yukawa couplings $h_{i\alpha} \lesssim (M_s/\mathrm{TeV})^{1/2} \times 10^{-3}$, so that the decay time of the right-handed neutrinos $\Gamma_{N_i}^{-1} \gtrsim (\mathrm{TeV}/M_s) \times 10^3 M_s^{-1}$. That is, they decay after the flaton has reached its minimum [but still long before the flaton decays, see Eq. (10)]. With $f_{ij} \sim 1$ in Eq. (1), the mass of the right-handed neutrinos when they decay is $M_{N_i} \sim M_I$. (The assumption $f_{ij} \sim 1$ can be relaxed without changing the main conclusions of this paper. For instance, we could have the third family right-handed neutrino mass $M_{N_3} \sim M_I$, whereas the first two family neutrinos are lighter.)

To be able to generate a lepton asymmetry, we must ensure that the right-handed neutrinos do not annihilate before they have time to decay. The annihilation rate for N_i through $B-L$ gauge interactions is $\Gamma_a \sim \Sigma_j |f_{ij}|^2 T^3 / 8\pi \langle \phi \rangle^2$. For $T \simeq T_c$, we estimate the annihilation probability P_a before the flaton reaches the minimum to be $P_a = 1 - \exp[\int_0^{\delta t} \Gamma_a \mathrm{d}t] \simeq \Sigma_j |f_{ij}|^2 / (120 \sigma^{1/2}) \sim 1/50$. Hence the number densities of N_i do not change significantly before they decay, at least not from this process. Similarly, the annihilation of N_i via dimension five couplings is also negligible. We therefore conclude that the N_i do not annihilate before they have time to decay.

The initial lepton asymmetry created by the decay of N_i is diluted by entropy production and also partially converted to the baryon asymmetry [9] by the sphaleron transition [10]. From the observed baryon asymmetry $n_B/s \simeq 8.7 \times 10^{-11}$ [16], the final lepton asymmetry is required to be $n_L/s \simeq 2.4 \times 10^{-10}$. To see how much initial lepton asymmetry is needed to account for this value, we first estimate the final temperature T_f and the dilution factor Δ.

The flaton, with mass $m_\phi = 2\sqrt{2} M_s$ mainly decays via the superpotential coupling $(\beta/M_*) \phi \bar{\phi} H_u H_d$. (Recall that the μ parameter, also induced by this term in the superpotential, is naturally of order M_s [$\mu \sim \beta M_I^2 / M_* \sim (\beta/\lambda) M_s$]). The decay width of the flaton is

$$\Gamma_\phi \simeq \frac{\beta^2}{8\pi} \frac{M_I^2}{M_*^2} m_\phi = \frac{1}{24\sqrt{2}\pi} \frac{\beta^2}{\lambda^2} \frac{M_s^3}{M_I^2}, \tag{10}$$

so that $\tau_\phi = \Gamma_\phi^{-1} \sim (M_I/M_s)^2 (\lambda^2/\beta^2) M_s^{-1} \gg M_s^{-1}$. For the final temperature $T_f \simeq 0.3 (\Gamma_\phi M_P)^{1/2}$, we find

$$T_f \simeq \left[\frac{\beta}{\lambda} \left(\frac{M_s}{\mathrm{TeV}} \right)^{3/2} \left(\frac{10^8 \ \mathrm{GeV}}{M_I} \right) \right] \times 15 \ \mathrm{TeV} \sim M_s \ \text{for} \ \beta \sim 0.1. \tag{11}$$

Note that the flaton decay products acquire a plasma mass $\sim gT$ [17] where g is the $B-L$ gauge coupling. The flaton decay thus can only take place once the temperature drops below $\sim m_\phi/g$. Consequently the final temperature T_f remains below $\sim m_\phi/g$ even for $M_I \ll 10^8$ GeV. We will assume for simplicity that the final temperature $T_f \sim M_s$.

It is gratifying that T_f is in a range where the electroweak sphalerons are able to convert some fraction of the lepton asymmetry into baryon asymmetry. This could not have been accomplished without the non-renormalizable term proportional to β in Eq. (2). [Integrating out N from Eq. (1) yields an effective dimension six operator which gives a ϕ decay rate $\Gamma \sim M_s^5 / M_I^4$, and the final temperature with only this decay would be of order GeV.] With $M_s \sim$ a few TeV, $\beta \sim 0.1$ leads to a μ term in the range of a few hundred GeV, as desired.

The entropy production due to ϕ decay dilutes the initial lepton asymmetry by a factor

$$\Delta \simeq \frac{4\mu_0^4 / 3T_f}{(2\pi^2/45) g_*(T_c) T_c^3} \simeq \frac{3\mu_0^4}{g_* T_c^3 T_f}, \tag{12}$$

where $g_* = 228.75$ for MSSM. Expressing the false vacuum energy density μ_0^4 and the critical temperature T_c in terms of M_s and M_I we obtain

$$\Delta \simeq 5 \times 10^{-3} \frac{\sigma^{3/2} M_I^2}{M_s T_f}. \tag{13}$$

We should make sure that the lepton asymmetry generated initially is large enough to sustain the impact of Δ. The lepton asymmetry after dilution is given by

$$\frac{n_L}{s} = \sum_i \frac{n_{N_i}}{s} \frac{1}{\Delta} \epsilon_i. \tag{14}$$

Here ϵ_i is the lepton asymmetry produced per decay of the i'th family neutrino N_i. Using Eqs. (6), (13), and $T_f \sim M_s$ we get

$$\frac{n_L}{s} \sim \sum_i 5 \left(\frac{0.2}{\sigma} \right)^{3/2} \left(\frac{M_s}{M_I} \right)^2 \epsilon_i. \tag{15}$$

For nearly degenerate neutrinos ϵ_i is given by [11,12]

$$\epsilon_i \simeq \sum_{j \neq i} \frac{\mathrm{Im}(h_{i\alpha}^* h_{\alpha j})^2}{|h_{i\alpha}|^2 |h_{j\alpha}|^2} \frac{\Delta M_N^2 M_{N_i} \Gamma_{N_j}}{(\Delta M_N^2)^2 + M_{N_i}^2 \Gamma_{N_j}^2}, \tag{16}$$

where $\Delta M_N^2 = M_{N_i}^2 - M_{N_j}^2 \simeq 2 M_{N_i} (M_{N_i} - M_{N_j})$. Defining $\xi_{ij} = (M_{N_i} - M_{N_j}) / (\Gamma_{N_j}/2)$, we can rewrite Eq. (16) as

$$\epsilon_i \simeq \sum_{j \neq i} \frac{\mathrm{Im}(h_{i\alpha}^* h_{\alpha j})^2}{|h_{i\alpha}|^2 |h_{j\alpha}|^2} \frac{\xi_{ij}}{\xi_{ij}^2 + 1}. \tag{17}$$

Assuming that $\delta_{CP} \equiv \mathrm{Im}(h_{i\alpha}^* h_{\alpha j})^2 / (|h_{i\alpha}|^2 |h_{j\alpha}|^2)$ is of order unity, the final lepton asymmetry is given by

$$\frac{n_L}{s}\sim\sum_{i,j\neq i} 2.4\times10^{-10}\left(\frac{0.2}{\sigma}\right)^{3/2}\left(\frac{M_s}{\mathrm{TeV}}\right)^2\left(\frac{10^8\mathrm{GeV}}{M_I}\right)^2\frac{\xi_{ij}}{\xi_{ij}^2+1}. \tag{18}$$

The asymmetry is maximized when the resonance condition $\xi_{ij}=1$ is satisfied for at least one pair of families. This gives an upper bound on the intermediate scale $M_I\sim(M_s/\mathrm{TeV})\times10^8$ GeV, which corresponds to a cutoff scale $M_*\sim(M_s/\mathrm{TeV})\times10^{14}$ GeV for $\lambda\sim1$. With somewhat larger values for M_s, say about 10 TeV, the upper bounds are $M_I\sim10^9$ GeV and $M_*\sim10^{15}$ GeV. One could ask how this relatively low cutoff scale can be incorporated within a more fundamental theory. One possibility is related to superstring inspired models with intermediate cutoff scales which have been of much recent interest. Another possibility is to introduce intermediate mass scale particles whose exchange can generate an effective cutoff scale of the desired magnitude, even though the underlying theory may have a cutoff scale that is significantly higher.

For $M_I\lesssim(M_s/\mathrm{TeV})\times10^8$ GeV the resonance condition does not have to be satisfied, although nearly degenerate right-handed neutrinos are still needed. Suppose that the neutrino mass differences $M_{N_i}-M_{N_j}$ are much greater than the decay widths $\Gamma_{N_j}(\xi_{ij}\gg1)$, so that $\epsilon_i\sim\Sigma_{j\neq i}\xi_{ij}^{-1}$. (Equation (17) in this case reduces to the perturbative result [18].) Using the seesaw relation $\Sigma_\alpha|h_{j\alpha}|^2\sim(m_{\nu_j}/\langle H_u\rangle^2)M_{N_j}$ with $M_{N_j}\sim f_jM_I$ (where f_j denotes an eigenvalue of f_{ij}), we can write

$$\Gamma_{N_j}=\sum_\alpha\frac{|h_{j\alpha}|^2M_{N_j}}{8\pi}\sim\frac{m_{\nu_j}f_j^2M_I^2}{8\pi\langle H_u\rangle^2}$$
$$\sim\left(\frac{m_{\nu_j}}{0.1\mathrm{eV}}\right)\left(\frac{M_I}{10^8\ \mathrm{GeV}}\right)^2 f_j^2\mathrm{GeV}. \tag{19}$$

Substituting Eq. (19) in Eq. (18), we obtain

$$\frac{n_L}{s}\sim\sum_{i,j\neq i} 2.4\times10^{-10}\left(\frac{0.2}{\sigma}\right)^{3/2}\left(\frac{M_s}{\mathrm{TeV}}\right)^2\left(\frac{m_{\nu_j}}{0.1\ \ \mathrm{eV}}\right)$$
$$\times\left(\frac{f_j^2\mathrm{GeV}}{M_{N_i}-M_{N_j}}\right). \tag{20}$$

The final lepton asymmetry is thus consistent with the observed baryon asymmetry provided that $M_s\ll M_I\lesssim(M_s/\mathrm{TeV})\times10^8$ GeV and that at least one pair of right-handed neutrino families have a mass difference less than or of order a GeV.

In conclusion, following Refs. [2,3], we have shown that in what are often referred to as thermal inflation models, there exists a novel mechanism for explaining the observed baryon asymmetry via leptogenesis. Because of significant entropy production that follows thermal inflation, the lepton asymmetry initially produced by heavy right-handed neutrinos with masses less than or of order M_I (but greater than the flaton mass) must be as large as possible. This requires nearly degenerate right-handed neutrinos with GeV scale mass differences. It remains to be seen how this degeneracy can be realized in conjunction with realistic neutrino masses and mixings. To ensure that the electroweak sphalerons can partially convert the lepton asymmetry to the observed baryon asymmetry, we require that the final temperature after completion of the phase transition is of order 10^3 GeV. This leads to the introduction of a term in the superpotential [Eq. (2)] which is also key to the resolution of the MSSM μ problem. Finally, it is clear that for intermediate scales significantly above 10^9 GeV, leptogenesis should arise from the decay products of the flaton field. For baryogenesis this has been discussed in Ref. [3].

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